

Gluino-squark contributions to CP violations in the kaon system

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Abstract

Recently it was shown, within the mass insertion approximation (MIA), that the gluino mediated flavor changing neutral current (FCNC) interactions can saturate both ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$ even for the real Cabibbo–Kobayashi–Maskawa (CKM) matrix through a single CP violating parameter $(\delta_{12}^d)_{LL}$. In this work, we extend our previous analysis to the nonvanishing KM phase, and to the effective SUSY model where the MIA is no longer a good approximation. In our model, ϵ_K and $B^0 - \overline{B}^0$ mixing can receive significant SUSY contributions. Therefore the usual constraints on ρ and η in the Wolfenstein parametrization of the CKM matrix can be relaxed except for the $|V_{ub}/V_{cb}|$ from the semileptonic $b \rightarrow u$ transition. This affects $K \rightarrow \pi\nu\bar{\nu}$ and $K_L \rightarrow \mu^+\mu^-$ indirectly, even if they are not directly induced by gluino-mediated FCNC, while $K_L \rightarrow \pi^0 e^+ e^-$ is influenced both directly and indirectly. Differences between our model and other SUSY models enhancing $\text{Re}(\epsilon'/\epsilon_K)$ are discussed in detail.

I. INTRODUCTION

For many years, CP violating phenomena was observed only in $K_L \rightarrow 2\pi$ [1], which could be attributed to $K^0 - \overline{K}^0$ mixing ($\Delta S = 2$). The CP violating parameter in the $K^0 - \overline{K}^0$ mixing, ϵ_K , has been accurately measured : $\epsilon_K = e^{i\pi/4} (2.271 \pm 0.017) \times 10^{-3}$ [2]. However, the recent observation of $\text{Re}(\epsilon'/\epsilon_K)$ by KTeV collaboration, $\text{Re}(\epsilon'/\epsilon_K) = (28.0 \pm 4.1) \times 10^{-4}$ [3], nicely confirms the earlier NA31 experiment [4] $\text{Re}(\epsilon'/\epsilon_K) = (23 \pm 7) \times 10^{-4}$. Including another new data from NA48 collaboration, $(14.0 \pm 4.3) \times 10^{-4}$ [5], the current world average for $\text{Re}(\epsilon'/\epsilon_K)$ is

$$\text{Re}(\epsilon'/\epsilon_K) = (19.2 \pm 2.4) \times 10^{-4}. \quad (1)$$

This nonvanishing number indicates unambiguously the existence of CP violation in the decay amplitude ($\Delta S = 1$), and the original form of superweak model is excluded. Along with the measurements of $\sin 2\beta$ from time-dependent asymmetry in $B^0 \rightarrow J/\psi K_S$ at B factories [6], we are now in the blooming era for detailed studies of CP violations in K and B systems. However, in this article, we will consider CP violations in the K system only for the reason to be explained later.

Two parameters ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$ that quantify CP violations in the kaon system can be accommodated by the Kobayashi–Maskawa (KM) phase in the Glashow–Salam–Weinberg’s standard model (SM). However, the SM predictions for the latter is rather uncertain because of large uncertainties in the nonperturbative matrix elements and the strange quark mass. There are significant variations in theoretical predictions [7–11], which indicate the current status of our understanding of $\text{Re}(\epsilon'/\epsilon_K)$. In this work, we will follow closely the work by Bosch *et al.* [7] in calculating $\text{Re}(\epsilon'/\epsilon_K)$ and other rare kaon decays so that we can compare the numerical results with other results more easily.

Although the SM can accommodate a bulk of the observed $\text{Re}(\epsilon'/\epsilon_K)$ within certain theoretical uncertainties, it is also interesting to speculate possible new physics contributions to CP violations in the kaon system. Among various scenarios beyond the SM, the minimal supersymmetric standard model (MSSM) and its various extensions are most promising candidates. Thus it is quite natural to consider generic SUSY contributions to ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$. In this regard, two of us (SB and PK) showed that the flavor preserving A_t and μ phases in the effective SUSY models alone cannot generate large enough ϵ_K [12–14] or new phase shift in $B^0 - \overline{B}^0$ mixing beyond that coming from the KM phase [15]. For ϵ'/ϵ_K , there have been a lot of works done after the KTeV announcement [16–29]. Those works in the supersymmetric models can be divided into three categories :

- the enhanced sdg vertex [19–27]
- the enhanced sdZ vertex [26,29]
- the enhanced EW penguins [28].

Models in the first category are divided further into two parts, depending on the origin of the flavor and CP violations : either from the flavor changing squark masses in the LL sector or the flavor changing LR sector due to nontrivial flavor structures in the trilinear A terms which may arise naturally in string inspired models with D branes. Our model presented

in Ref. [23] is unique compared to other models with enhanced sdg vertex, since SUSY CP violations in our model affects not only the $\text{Re}(\epsilon'/\epsilon_K)$ but also the ϵ_K by significant amounts, contrary to other models in this category. In fact, it was shown that both ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$ can be saturated by a single complex number $(\delta_{12}^d)_{LL}$, if $|\mu \tan \beta|$ is large enough $\sim \mathcal{O}(10)$ TeV so that the induced $(\delta_{12}^d)_{LR}$ through the double mass insertion can be $\mathcal{O}(10^{-5})$. Furthermore, there is no conflict with the neutron electric dipole moment (EDM) in our model. There is another work showing that all the CP violations observed so far can be accommodated by a single CP violating phase in the gluino mass parameter $M_{\tilde{g}}$ [24]. They assumed a very specific flavor structure in the trilinear A couplings in order to generate large ϵ' . On the other hand, we assume a specific flavor structure in the left squark mass matrix only, and the flavor structure and CP violating phase of a trilinear A coupling is totally irrelevant in our case.

In generic SUSY models, there can be potentially large contributions to flavor changing neutral currents (FCNC) and CP violating processes from gluino-squark mediations, which is termed as SUSY flavor/CP problems [30]. Since these effects are due to strong interactions, it may be parametrically large compared to the charged Higgs and chargino contributions, unless the squarks and gluinos are too heavy. There are basically three ways to solve SUSY flavor problems : (i) universality of squark masses, (ii) alignment of squark and quark mass matrices due to some flavor symmetries, and (iii) decoupling of the 1st/2nd generation squarks.

In order to study the gluino mediated flavor changing phenomena within the frameworks of approximate universality or alignment, it is convenient to use the so-called mass insertion approximation (MIA) [31]. In this approximation, one works in the super CKM basis where the quark mass matrices are diagonal. The quark-squark-gluino vertex is flavor diagonal in the MIA, and the flavor/chirality mixing occur through the insertion of $(\delta_{ij}^d)_{AB}$, where $i, j = 1, 2, 3$ and $A, B = L, R$ denote the flavors of the squarks under consideration and the chiralities of their superpartners. The superscript denotes that the down type squark mass matrix is involved. The parameters $(\delta_{ij}^d)_{AB}$ characterize the size of the gluino-mediated flavor changing amplitudes, and they may be CP violating complex numbers, in general. In the following, diagrams involving charged Higgs, chargino and neutralino will be ignored, since they are suppressed by α_w/α_s compare to the gluino-squark loops unless gluino/squarks are very heavy. This should be a good starting point for studying the SUSY FCNC/CP problems.

In addition to the MIA, we also consider the vertex mixing (VM) model which is a good approximation if the third generation sfermions are lighter than the 1st/2nd generation sfermions as in the effective SUSY models [32]. In this case, the SUSY FCNC and CP problems are solved by the decoupling (and some degeneracy in the 1st/2nd generation sfermion masses), and we keep only the left/right sbottoms and gluino contributions to the $\Delta S = 1, 2$ effective Hamiltonians. The related works in B physics can be found in Refs. [33].

Although SUSY CP violations can saturate both ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$, it would be unnatural to assume that δ_{KM} vanishes (or equals π), since there is no symmetry principle supporting such assumption ¹. In this work, we present more detailed analysis of our model within both

¹We can impose CP symmetry to be broken only softly. This case is already included in our model

MIA and VM cases with nonvanishing δ_{KM} . Since the ϵ_K and $B^0 - \overline{B}^0$ mixing are affected by SUSY contributions significantly in our model, the usual constraints on (ρ, η) from these two quantities need not be held any more. Only the constraint from $b \rightarrow u$ semileptonic decay would be valid. Therefore, when we vary the (ρ, η) , we will impose this constraint only, and the apex of the unitarity triangle (UT) can be anywhere in the (ρ, η) plane, not only in the first quadrangle. This is in sharp contrast to other analysis within SUSY models where new physics contributes significantly only to $\text{Re}(\epsilon'/\epsilon_K)$ or the analysis based on the MSSM with minimal flavor violation. There is an indication that the unitarity triangle does not collapse (namely α, β, γ do not vanish) just from determining the triangle only from the unitarity of the Cabibbo–Kobayashi–Maskawa (CKM) matrix [35] :

$$\alpha = 19^\circ - 142^\circ, \quad \beta = 6^\circ - 31^\circ, \quad \gamma = 28^\circ - 152^\circ. \quad (2)$$

Also there are some indications that charmless nonleptonic 2-body decays of B mesons seem to prefer $\cos \gamma < 0$ [36,37], although theoretical uncertainties involved are of vastly different degrees. Considering all these points, we consider all the possibility, although $\gamma \sim \pi$ may be unrealistic.

Assuming that the ϵ_K and $B^0 - \overline{B}^0$ mixing cannot be used to constrain the CKM elements, we study several rare K decays such as $K \rightarrow \pi \nu \bar{\nu}$, $K \rightarrow \pi l^+ l^-$ and $K \rightarrow \mu^+ \mu^-$ including the squark–gluino contributions, and compared with other recent works on $\text{Re}(\epsilon'/\epsilon_K)$. We do not consider the CP asymmetries in $K \rightarrow 3\pi$, $K \rightarrow \pi\pi\gamma$ and hyperon decays. These observables were discussed in the literatures within the context of the enhanced sdg vertex as the origin of large ϵ'/ϵ_K . In our model, we have another source of FCNC and CP violation, namely $(\delta_{12}^d)_{LL}$ which induces $\Delta S = 1$ four-quark operators. The detailed study of these effects to $K \rightarrow 3\pi$, $K \rightarrow \pi\pi\gamma$ and hyperon decays are beyond the scope of the present work. Also these processes will suffer from large hadronic uncertainties as other nonleptonic kaon decays in chiral perturbation theory, unlike $K \rightarrow \pi \nu \bar{\nu}$.

The contents of this paper are organized as follows. In the Sec. II A and B, we review briefly the $\Delta S = 2$ and $\Delta S = 1$ effective Hamiltonians and their relations to ϵ_K and ϵ'/ϵ_K . In Sec. II C and D, we recapitulate the basic formulae for neutron EDM, and branching ratios and CP asymmetries for $K \rightarrow \pi \nu \bar{\nu}$, $K \rightarrow \pi l^+ l^-$ and $K \rightarrow \mu^+ \mu^-$. The numerical inputs for various parameters in our model are summarized in Sec. II E. In Sec. III and IV, we discuss the CP violation phenomenology in the kaon sector for the MIA and the VM cases, respectively. Finally we summarize in Sec. V.

II. EFFECTIVE HAMILTONIANS FOR KAON PHYSICS

In order to discuss ϵ_K , ϵ'/ϵ_K and other rare kaon decays within the SM and SUSY models, we first construct the effective Hamiltonians for $\Delta S = 2$ and $\Delta S = 1$. We obtain the Wilson coefficients of the effective theory after matching it to the full theory at $\mathcal{O}(m_W)$. In this section, we define the operator basis for these effective Hamiltonians and the recipes for the renormalization group (RG) running from the weak scale to the lower energy scale $\mu \approx m_c$.

with $\delta_{KM} = 0$ or π . See Refs. [34] for more discussions.

The relevant Wilson coefficients for the $\Delta S = 1$ and $\Delta S = 2$ effective Hamiltonians are explicitly given in the Secs. III and IV, where we describe the MI and the VM approximations for the CP and flavor-violating effects of SUSY in detail.

A. Effective Hamiltonian for $\Delta S = 2$ and ϵ_K

The effective Hamiltonian describing the $\Delta S = 2$ $K^0 - \overline{K}^0$ mixing in the SM and the MSSM (with only the gluino contributions included) can be written as

$$H_{\text{eff}}^{\Delta S=2} = H_{\text{SM}}^{\Delta S=2} + H_{\text{SUSY}}^{\Delta S=2} \quad (3)$$

where [38,39]

$$H_{\text{SM}}^{\Delta S=2} = \frac{G_F^2}{16\pi^2} m_W^2 \left[\lambda_c^{*2} \eta_1^{(SM)} S_0(x_c) + \lambda_t^{*2} \eta_2^{(SM)} S_0(x_t) + 2\lambda_c^* \lambda_t^* \eta_3^{(SM)} S_0(x_c, x_t) \right] \\ \times [\alpha_s(\mu)]^{-2/9} \bar{d}^\alpha \gamma_\mu (1 - \gamma_5) s^\alpha \bar{d}^\beta \gamma^\mu (1 - \gamma_5) s^\beta + \text{h.c.}, \quad (4)$$

$$H_{\text{SUSY}}^{\Delta S=2} = \frac{1}{216 \widetilde{m}^2} \left(\sum_{i=1}^5 C_i^{(2)} Q_i + \sum_{i=1}^3 \tilde{C}_i^{(2)} \tilde{Q}_i \right), \quad (5)$$

with operators defined to be

$$\begin{aligned} Q_1 &= \alpha_s^2(\mu) \bar{d}^\alpha \gamma_\mu (1 - \gamma_5) s^\alpha \bar{d}^\beta \gamma^\mu (1 - \gamma_5) s^\beta, \\ Q_2 &= \alpha_s^2(\mu) \bar{d}^\alpha (1 - \gamma_5) s^\alpha \bar{d}^\beta (1 - \gamma_5) s^\beta, \\ Q_3 &= \alpha_s^2(\mu) \bar{d}^\alpha (1 - \gamma_5) s^\beta \bar{d}^\beta (1 - \gamma_5) s^\alpha, \\ Q_4 &= \alpha_s^2(\mu) \bar{d}^\alpha (1 - \gamma_5) s^\alpha \bar{d}^\beta (1 + \gamma_5) s^\beta, \\ Q_5 &= \alpha_s^2(\mu) \bar{d}^\alpha (1 - \gamma_5) s^\beta \bar{d}^\beta (1 + \gamma_5) s^\alpha. \end{aligned} \quad (6)$$

Here, \widetilde{m} is the average squark mass in MIA, or the sbottom mass in VM. α and β are color indices. The operators \tilde{Q}_i are obtained from Q_i by interchanging L and R . The explicit form of the loop function, $S_0(x, y)$, can be found in Ref. [38].

We have separated the effective Hamiltonian into two parts : one coming from the SM and the other from SUSY sector, because the former is of $\mathcal{O}(\alpha_w)$, whereas the latter is of $\mathcal{O}(\alpha_s)$. We separately do RG running of SUSY and SM contributions to get the Wilson coefficients for the kaon physics at $\mathcal{O}(m_c)$. In the SM part, we use the well-known anomalous dimension matrices for RG running which can be found in Ref. [38]. For the RG running of SUSY part, we use a modified method advocated in [40] in which new effective operators are defined including the α_s^n factor coming from the gluino-squark loop diagrams.

After defining the new operator set, we can apply the RG procedure of SM to gluino mediated contributions: *matching* calculation for Wilson coefficients to order α_s^0 and *RG evolution* by using the anomalous dimension matrix to order α_s^1 for leading log approximations (LLA). The calculation of matrix elements is obviously modified by the additional factors of new operators.

The anomalous dimension matrix in our basis at leading order is given by

$$\gamma^{(0)} = \tilde{\gamma}^{(0)} + 4 \beta_0^{(f)} \mathbf{1}, \quad (7)$$

where $\beta_0^{(f)} = (11 - 2f/3)$ with the effective flavor number of f and $\tilde{\gamma}^{(0)}$ can be obtained in [41],

$$\tilde{\gamma}^{(0)} = \begin{pmatrix} 4 & 0 & 0 & 0 & 0 \\ 0 & -28/3 & 4/3 & 0 & 0 \\ 0 & 16/3 & 23/3 & 0 & 0 \\ 0 & 0 & 0 & -16 & 0 \\ 0 & 0 & 0 & -6 & 0 \end{pmatrix}. \quad (8)$$

The Wilson coefficients at $\mathcal{O}(m_c)$ running down from those at $\mathcal{O}(m_W)$ read

$$\begin{aligned} C_1^{(2)}(m_c) &= \eta_1 C_1^{(2)}(m_W), \\ C_2^{(2)}(m_c) &= \eta_{22} C_2^{(2)}(m_W) + \eta_{23} C_3^{(2)}(m_W), \\ C_3^{(2)}(m_c) &= \eta_{32} C_2^{(2)}(m_W) + \eta_{33} C_3^{(2)}(m_W), \\ C_4^{(2)}(m_c) &= \eta_4 C_4^{(2)}(m_W) + \frac{1}{3} (\eta_4 - \eta_5) C_5^{(2)}(m_W), \\ C_5^{(2)}(m_c) &= \eta_5 C_5^{(2)}(m_W), \end{aligned} \quad (9)$$

with

$$\begin{aligned} \eta_1 &= \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{56/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{52/23}, \\ \eta_2 &= \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{1.419} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{1.369}, \\ \eta_3 &= \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{2.661} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{2.718}, \\ \eta_4 &= \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{26/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{22/23}, \\ \eta_5 &= \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{53/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{49/23}, \end{aligned} \quad (10)$$

and

$$\begin{aligned} \eta_{22} &= 0.983\eta_2 + 0.017\eta_3, & \eta_{23} &= -0.258\eta_2 + 0.258\eta_3, \\ \eta_{32} &= -0.064\eta_2 + 0.064\eta_3, & \eta_{33} &= 0.017\eta_2 + 0.983\eta_3. \end{aligned} \quad (11)$$

In order to estimate ΔM_K and ϵ_K , it is necessary to know the hadronic matrix elements: $\langle \bar{K}^0 | Q_i(\mu) | K^0 \rangle$ for $\mu \sim m_c$. In the present work, we use the following expressions [42,43], which are sufficient for our purpose in the LLA.

$$\begin{aligned} \langle K^0 | Q_1(\mu) | \bar{K}^0 \rangle &= \frac{8}{3} \alpha_s^2(\mu) m_K^2 f_K^2 B_1(\mu), \\ \langle K^0 | Q_2(\mu) | \bar{K}^0 \rangle &= -\frac{5}{3} \alpha_s^2(\mu) \left(\frac{m_K}{m_s(\mu) + m_d(\mu)} \right)^2 m_K^2 f_K^2 B_2(\mu), \end{aligned}$$

$$\begin{aligned}
\langle K^0 | Q_3(\mu) | \overline{K^0} \rangle &= \frac{1}{3} \alpha_s^2(\mu) \left(\frac{m_K}{m_s(\mu) + m_d(\mu)} \right)^2 m_K^2 f_K^2 B_3(\mu), \\
\langle K^0 | Q_4(\mu) | \overline{K^0} \rangle &= \alpha_s^2(\mu) \left[\frac{1}{3} + 2 \left(\frac{m_K}{m_s(\mu) + m_d(\mu)} \right)^2 \right] m_K^2 f_K^2 B_4(\mu), \\
\langle K^0 | Q_5(\mu) | \overline{K^0} \rangle &= \alpha_s^2(\mu) \left[1 + \frac{2}{3} \left(\frac{m_K}{m_s(\mu) + m_d(\mu)} \right)^2 \right] m_K^2 f_K^2 B_5(\mu).
\end{aligned} \tag{12}$$

The bag parameters B 's are given by [43]

$$\begin{aligned}
B_1(\mu) &= 0.60, \\
B_2(\mu) &= 0.66, \\
B_3(\mu) &= 1.05, \\
B_4(\mu) &= 1.03, \\
B_5(\mu) &= 0.73,
\end{aligned} \tag{13}$$

at $\mu = 2 \text{ GeV}$ (We use these values at $\mu = m_c = 1.3 \text{ GeV}$).

The CP violation in the $K^0 - \overline{K^0}$ mixing is characterized by a complex parameter ϵ_K , which can be obtained from the $\Delta S = 2$ effective Hamiltonian by the following formula:

$$\epsilon_K = \frac{\exp(i\pi/4)}{\sqrt{2}\Delta M_K} [\text{Im}M_{12} + 2\xi \text{Re}M_{12}], \tag{14}$$

where M_{12} is defined as

$$2m_K M_{12} = \langle K^0 | H_{\text{eff}}^{\Delta S=2} | \overline{K^0} \rangle. \tag{15}$$

Remember that $\xi = \text{Im}A_0/\text{Re}A_0$ is very small and thus is neglected in our calculation.

B. Effective Hamiltonian for $\Delta S = 1$ and $\text{Re}(\epsilon'/\epsilon_K)$

The effective Hamiltonian describing the $\Delta S = 1$ kaon decays receive contributions from the SM and reads (with only the gluino contributions included):

$$H_{\text{eff}}^{\Delta S=1} = H_{\text{SM}}^{\Delta S=1} + H_{\text{SUSY}}^{\Delta S=1}, \tag{16}$$

where

$$H_{\text{SM}}^{\Delta S=1} = C'_1 O'_1 + C'_2 O'_2 + \sum_{i=3}^{10} C'_i O'_i + \text{h.c.}, \tag{17}$$

$$H_{\text{SUSY}}^{\Delta S=1} = C_1 O_1 + C_2 O_2 + \sum_{i=3}^{10} (C_i O_i + \tilde{C}_i \tilde{O}_i) + \sum_{A=\pm} \sum_{B=\gamma, g} \sum_{X=s, \tilde{g}} C_{BX}^A O_{BX}^A + \text{h.c.} \tag{18}$$

The explicit form of O'_i can be referred in Ref. [38] and new operators O_i 's are defined as

$$O_i(\mu) = \alpha_s^2(\mu) O'_i(\mu). \tag{19}$$

\tilde{O}_i are obtained from O_i by interchanging $1 - \gamma_5$ and $1 + \gamma_5$. The 10×10 anomalous dimension matrix for RG evolution in these new four-quark operators is defined in a similar way to that of $\Delta S = 2$ case given in Eq. (7).

The last part of Eq.(18) is for the magnetic and chromomagnetic operators in which the subindex B denotes whether the chirality flips on the strange quark or the gluino line:

$$\begin{aligned} O_{\gamma s}^\pm &= m_s \frac{e \alpha_s(\mu)}{32\pi^2} \left[\bar{d} \sigma^{\mu\nu} F_{\mu\nu} (1 + \gamma_5) s \pm \bar{d} \sigma^{\mu\nu} F_{\mu\nu} (1 - \gamma_5) s \right], \\ O_{\gamma \tilde{g}}^\pm &= \frac{e \alpha_s(\mu)}{32\pi^2} \left[\bar{d} \sigma^{\mu\nu} F_{\mu\nu} (1 + \gamma_5) s \pm \bar{d} \sigma^{\mu\nu} F_{\mu\nu} (1 - \gamma_5) s \right], \\ O_{gs}^\pm &= m_s \frac{g \alpha_s(\mu)}{32\pi^2} \left[\bar{d} \sigma^{\mu\nu} t^a G_{\mu\nu}^a (1 + \gamma_5) s \pm \bar{d} \sigma^{\mu\nu} t^a G_{\mu\nu}^a (1 - \gamma_5) s \right], \\ O_{g\tilde{g}}^\pm &= \frac{g \alpha_s(\mu)}{32\pi^2} \left[\bar{d} \sigma^{\mu\nu} t^a G_{\mu\nu}^a (1 + \gamma_5) s \pm \bar{d} \sigma^{\mu\nu} t^a G_{\mu\nu}^a (1 - \gamma_5) s \right]. \end{aligned}$$

The anomalous dimension matrices for RG evolution of the above equations are given by [40]

$$\gamma_{\tilde{g}}^{(0)} = \tilde{\gamma}^{(0)} + (2\beta_0^{(f)} - 8) \mathbf{1}, \quad (20)$$

$$\gamma_s^{(0)} = \tilde{\gamma}^{(0)} + 2\beta_0^{(f)} \mathbf{1}, \quad (21)$$

$$\tilde{\gamma}^{(0)} = \begin{pmatrix} 32/3 & 0 \\ -32/9 & 28/3 \end{pmatrix}. \quad (22)$$

The Wilson coefficients of magnetic operators at $\mathcal{O}(m_c)$ are

$$\begin{aligned} C_{\gamma(\tilde{g},s)}(m_c) &= \eta_1^{(\tilde{g},s)} C_{\gamma(\tilde{g},s)}(m_W) - \frac{8}{3} \left(\eta_1^{(\tilde{g},s)} - \eta_2^{(\tilde{g},s)} \right) C_{g(\tilde{g},s)}(m_W), \\ C_{g(\tilde{g},s)}(m_c) &= \eta_2^{(\tilde{g},s)} C_{g(\tilde{g},s)}(m_W), \end{aligned} \quad (23)$$

with

$$\eta_1^{\tilde{g}} = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{29/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{27/23}, \quad \eta_2^{\tilde{g}} = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{27/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{25/23}, \quad (24)$$

and

$$\eta_1^s = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{41/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{39/23}, \quad \eta_2^s = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)} \right)^{39/25} \left(\frac{\alpha_s(m_W)}{\alpha_s(m_b)} \right)^{37/23}. \quad (25)$$

The hadronic matrix element of a four quark operator in Eq. (19) is obtained by multiplying the corresponding one of SM in Ref. [38] by $\alpha_s^2(m_c)$. For the magnetic operators, the matrix elements at $\mathcal{O}(m_c)$ can be expressed as [26]

$$\langle O_{gs}^- \rangle_0 = m_s \langle O_{g\tilde{g}}^- \rangle_0 = m_s \alpha_s(\mu) \sqrt{\frac{3}{2}} \frac{11}{16\pi^2} \frac{\langle \bar{q}q \rangle}{F_\pi^3} m_\pi^2 B_G, \quad \langle O_g^+ \rangle_0 = 0, \quad (26)$$

where the subscript 0 means $\Delta I = 0$ components of the matrix elements.

The parameter $\text{Re}(\epsilon'/\epsilon_K)$, which characterizes the direct CP violation in the decay amplitude of $K \rightarrow \pi\pi$, is expressed as following in terms of the $\Delta S = 1$ effective Hamiltonian:

$$\text{Re}\left(\frac{\epsilon'}{\epsilon_K}\right) = -\frac{\omega}{\sqrt{2}|\epsilon|\text{Re}A_0} \sum_i \text{Im}(C_i) \left(\langle O_i \rangle_0 (1 - \Omega_{\eta+\eta'}) - \frac{1}{\omega} \langle O_i \rangle_2 \right) \quad (27)$$

where $\omega = \text{Re}A_2/\text{Re}A_0 = 0.045$, and $\Omega_{\eta+\eta'} = 0.25 \pm 0.05$ represents the isospin breaking effects.

C. Effective Hamiltonian for the neutron EDM

The electric dipole moment of a fermion f can be described by the effective Hamiltonian [44]

$$H_{\text{eff}}^{\text{EDM}} = \sum_{i=1}^3 C_i^{\text{edm}} O_i^{\text{edm}}, \quad (28)$$

and its basis operators

$$\begin{aligned} O_1^{\text{edm}} &= -\frac{i}{2} \bar{f} \sigma^{\mu\nu} \gamma_5 f F_{\mu\nu}, \\ O_2^{\text{edm}} &= -\frac{i}{2} \bar{f} \sigma^{\mu\nu} \gamma_5 T^a f G_{\mu\nu}^a, \\ O_3^{\text{edm}} &= -\frac{1}{6} f_{abc} G_{\mu\rho}^a G_{\nu}^{b\rho} G_{\lambda\sigma}^c \epsilon^{\mu\nu\lambda\sigma}. \end{aligned} \quad (29)$$

With this Hamiltonian, we calculate the (chromo)EDM of quarks and gluons and compare the result to the experimental value to constrain possible parameter space. The current experimental upper bound on the neutron EDM is [45]

$$|d_n| \lesssim 6.3 \times 10^{-26} e \text{ cm}. \quad (30)$$

D. Rare kaon decays : $K \rightarrow \pi\nu\bar{\nu}$, $\pi l^+ l^-$, $\mu^+ \mu^-$

SUSY contributions to various rare kaon decays through the (chromo)magnetic and enhanced sdZ operators have been discussed comprehensively by Buras *et al.* [26], and we heavily lean on formulae for branching ratios for rare kaon decays given in their paper. In the numerical analysis, however, there are two major different points to be emphasized :

- In our case, we ignore the enhanced sdZ operator, which may be important only in a restricted parameter space, as already noticed in Ref. [26]. Therefore, our model is precisely the Scenario A therein.
- When we vary CKM matrix elements λ_t , we impose constraints coming from semileptonic decays only, since these are fairly insensitive to new physics beyond the SM, especially on the SUSY contributions we are considering here. On the other hand,

ϵ_K and $B_{d(s)}^0 - \overline{B_{d(s)}^0}$ mixing may be affected by new physics. Indeed gluino mediated FCNC present in our model can affect both of these observables in a significant manner, unlike those considered in most recent works related with ϵ'/ϵ_K in the framework of SUSY. Since we consider the kaon sector only and the B meson sector is independent of the kaon sector in the MIA, we do not impose the $B^0 - \overline{B^0}$ mixing in this work. Therefore, the weak phase $\gamma = \phi_3$ can be anywhere between 0 and 2π .

Also $(\delta_{12}^d)_{LL}$ would contribute to $s \rightarrow dq\bar{q}$ penguin operators, thereby modifying nonleptonic kaon decays (and CP violations therein) such as $K \rightarrow 3\pi$ and $K \rightarrow \pi\pi\gamma$ and hyperon decays. The effects of $(\delta_{12}^d)_{LR}$ on these decays have been considered in the literatures, but not those of $(\delta_{12}^d)_{LL}$. Since these will involve some hadronic uncertainties compared to rare kaon decays we consider here (such as $K \rightarrow \pi\nu\bar{\nu}$), we will not include these nonleptonic kaon and hyperon decays in this work.

In the SM, $K \rightarrow \pi\nu\bar{\nu}$ is affected by the KM phase. In fact the neutral mode $K_L \rightarrow \pi^0\nu\bar{\nu}$ is purely CP violating so that its branching ratio would vanish in the limit of exact CP symmetry. In our model with two sources of CP violations, only δ_{KM} affects these decays, since $s \rightarrow dg$ vertex does not contribute to them.

Using the formulae given in Ref. [26] we have the results

$$B(K_L \rightarrow \pi^0\nu\bar{\nu}) = B_{\text{SM}}^0 = 6.78 \times 10^{-4} (X_0 \text{Im}\lambda_t)^2, \quad (31)$$

$$B(K^+ \rightarrow \pi^+\nu\bar{\nu}) = B_{\text{SM}}^+ = 1.55 \times 10^{-4} \left[(X_0 \text{Im}\lambda_t)^2 + (X_0 \text{Re}\lambda_t + \Delta_c)^2 \right], \quad (32)$$

where $X_0 = C_0 - 4B_0 = 1.52$ is a combination of loop functions evaluated at $\bar{m}_t(m_t) = 166$ GeV, and $\Delta_c = -(2.11 \pm 0.30) \times 10^{-4}$ is the internal charm contribution.

Unlike $K \rightarrow \pi\nu\bar{\nu}$, the direct CP violating part of $K_L \rightarrow \pi^0 e^+ e^-$ is affected in our model, because the Wilson coefficients of $O_{\gamma(\tilde{g},s)}$ and $O_{7V} \equiv (\bar{d}s)_{(V-A)}(\bar{e}e)_V$ are modified due to the new CP violating phase from the squark–gluino diagrams. As for other SUSY contributions, we work with the ‘SUSY version’ of operator

$$O_{7V}^{\text{SUSY}'} \equiv \alpha \alpha_s(\mu) (\bar{d}s)_{(V-A)}(\bar{e}e)_V, \quad (33)$$

rather than O_{7V} , in order to do matching and running formally consistent with LLA. Under the definition,

$$\text{Im}\Lambda_g^+ \tilde{y}_\gamma = -\frac{\alpha_s B_T}{\sqrt{2} G_F m_K} \left[m_s \text{Im}C_{\gamma s}^+ + \text{Im}C_{\gamma \tilde{g}}^+ \right], \quad (34)$$

the branching ratio is given by

$$B(K_L \rightarrow \pi^0 e^+ e^-)_{\text{dir}} = 6.3 \times 10^{-6} \times \left[\left(\text{Im}\lambda_t \tilde{y}_{7V} + \text{Im}\Lambda_g^+ \tilde{y}_\gamma + \frac{\sqrt{2}}{G_F} 2\pi\alpha_s(m_c) \text{Im}C_{7V}^{\text{SUSY}'} \right)^2 + (\text{Im}\lambda_t \tilde{y}_{7A})^2 \right], \quad (35)$$

where, for $\Lambda_{\overline{MS}} = 325$ MeV and $m_t = 170$ GeV, [38]

$$\begin{aligned}\tilde{y}_{7V}(\mu = 1.3 \text{ GeV}) &= 0.537 \times 2\pi, \\ \tilde{y}_{7A}(\mu = 1.3 \text{ GeV}) &= -0.700 \times 2\pi.\end{aligned}$$

The short distance part of the branching ratio of $K_L \rightarrow \mu^+ \mu^-$ within our model is the same as that in the SM [26]:

$$B(K_L \rightarrow \mu^+ \mu^-)_{\text{SD}} = B_{\text{SM}}^{\mu\mu} = 6.32 \times 10^{-3} (Y_0 \text{Re}\lambda_t + \overline{\Delta}_c)^2, \quad (36)$$

where $\overline{\Delta}_c = -(6.54 \pm 0.60) \times 10^{-5}$ is the charm contribution.

E. Input parameters

Since there may be large SUSY contributions to ϵ_K as well as to $B^0 - \overline{B}^0$ in our model, one cannot use the constraints on the CKM matrix elements coming from these observables. We impose only the constraint coming from the semileptonic decays of mesons. Therefore we have [2]

$$|V_{ub}/V_{cb}| = 0.08 \pm 0.02 \quad (37)$$

This leaves the phase of V_{ub} undetermined so that the SM contributions to various observables are not fully determined. In the following numerical analysis, we sample data fixing the phase of V_{ub} to some selected values of $0^\circ, 60^\circ, 120^\circ \dots$ so on up to 360° .

We list numerical input values used throughout our analysis in Table I.

III. MASS INSERTION APPROXIMATION CASE

A. Approximate flavor symmetry and $(\delta_{12}^d)_{AB}$

One way to solve the gluino-mediated FCNC problems in general SUSY models is to assume that the squark masses are approximately diagonal and degenerate in the basis where quark mass matrices are diagonal (the so-called super CKM basis).

$$(\tilde{m}_{ij}^2)_{AB} = (\tilde{m}^2)_{AB} \delta_{ij} + (\delta m_{ij}^2)_{AB}.$$

In this case, one can expand the squark propagator around the common squark mass matrix $(\tilde{m}^2)_{AB} \delta_{ij}$, and the rest can be cast into interaction Hamiltonians. As long as δm^2 is small compared to the common squark mass \tilde{m}^2 , one can keep the lowest order perturbations in δm^2 relevant to the processes we are interested in. In many cases, just a single mass insertion along a single squark propagator is sufficient, but double mass insertion along a single squark propagator should be included in some cases. In fact, it is the case for ϵ'/ϵ_K in the relatively large $\mu \tan \beta$ region [23]. In this case, smallness of s quark mass can be compensated by a large $\mu \tan \beta$ (which has been ignored in most previous literatures) so that an appreciable amount of $\tilde{s}_R - \tilde{s}_L$ may be possible.

The typical size of $(\delta_{ij}^d)_{AB} \equiv (\delta m_{ij}^2)_{AB}/\tilde{m}^2$ should be fairly small in order to be consistent with ΔM_K and ϵ_K . Theoretical understanding of such a small number constitutes the so-called SUSY flavor problems. This problem could be solved by approximate universality or

approximate flavor symmetry where both quarks and squarks are almost aligned in flavor space.

The gluino–squark contributions to the $\Delta S = 1, 2$ effective Hamiltonian at the heavy SUSY particle scale have been calculated by several groups, and are encoded in the Wilson coefficients listed below.

B. Wilson coefficient for $\Delta S = 2$

The Wilson coefficients for the $\Delta S = 2$ effective Hamiltonian at the SUSY particle mass scales coming from the gluino contributions are given as follows:

$$\begin{aligned} C_1^{(2)} &= -\frac{1}{4} \left(24x f_6(x) + 66 \tilde{f}_6(x) \right) \left(\delta_{12}^d \right)_{LL}^2, \\ C_2^{(2)} &= C_3^{(2)} = C_4^{(2)} = C_5^{(2)} = \tilde{C}_1^{(2)} = 0, \\ \tilde{C}_2^{(2)} &= -\frac{816}{4} x f_8(x) \left(\delta_{12}^d \right)_{LL}^2 \left(\delta_{22}^d \right)_{LR}^2, \\ \tilde{C}_3^{(2)} &= \frac{144}{4} x f_8(x) \left(\delta_{12}^d \right)_{LL}^2 \left(\delta_{22}^d \right)_{LR}^2. \end{aligned} \quad (38)$$

The expressions for $\tilde{C}_2^{(2)}$ and $\tilde{C}_3^{(2)}$ are new, but their contributions to ϵ_K and Δm_K are negligible. The loop functions are given by the following expressions :

$$\begin{aligned} f_6(x) &= \frac{6(1+3x) \ln x + x^3 - 9x^2 - 9x + 17}{6(x-1)^5}, \\ \tilde{f}_6(x) &= \frac{6x(1+x) \ln x - x^3 - 9x^2 + 9x + 1}{3(x-1)^5}, \\ f_8(x) &= \frac{197 + 25x - 300x^2 + 100x^3 - 25x^4 + 3x^5 + 60(1+5x) \ln x}{60(x-1)^7}, \end{aligned} \quad (39)$$

where $x = (m_{\tilde{g}}/\tilde{m})^2$. The flavor conserving but chirality changing mass insertion parameter $(\delta_{22}^d)_{LR}$ is defined as

$$(\delta_{22}^d)_{LR} = \frac{m_s}{\tilde{m}^2} (A_s^* - \mu \tan \beta) \quad (40)$$

We neglect the terms proportional to $(\delta_{12}^d)_{LR}^2$, $(\delta_{12}^d)_{RL}^2$, and $(\delta_{12}^d)_{RR}^2$.

C. Wilson coefficient for $\Delta S = 1$

The Wilson coefficients for the $\Delta S = 1$ effective Hamiltonian at the SUSY particle mass scales are given as follows :

$$\begin{aligned} C_i &= 0 \quad (\text{for } i = 1, 2, 7-10) \\ C_3 &= \frac{1}{4\tilde{m}^2} \left(-\frac{1}{9} B_1(x) - \frac{5}{9} B_2(x) - \frac{1}{18} P_1(x) - \frac{1}{2} P_2(x) \right) \left(\delta_{12}^d \right)_{LL} \end{aligned}$$

$$\begin{aligned}
C_4 &= \frac{1}{4\widetilde{m}^2} \left(-\frac{7}{3}B_1(x) + \frac{1}{3}B_2(x) + \frac{1}{6}P_1(x) + \frac{3}{2}P_2(x) \right) (\delta_{12}^d)_{LL} \\
C_5 &= \frac{1}{4\widetilde{m}^2} \left(\frac{10}{9}B_1(x) + \frac{1}{18}B_2(x) - \frac{1}{18}P_1(x) - \frac{1}{2}P_2(x) \right) (\delta_{12}^d)_{LL} \\
C_6 &= \frac{1}{4\widetilde{m}^2} \left(-\frac{2}{3}B_1(x) + \frac{7}{6}B_2(x) + \frac{1}{6}P_1(x) + \frac{3}{2}P_2(x) \right) (\delta_{12}^d)_{LL} \\
\tilde{C}_i &= C_i \text{ with replacement of } LL \text{ by } RR, \quad i = 3, 4, 5, 6
\end{aligned} \tag{41}$$

where

$$\begin{aligned}
B_1(x) &= \frac{1 + 4x - 5x^2 + 4x\log x + 2x^2\log x}{8(1-x)^4} \\
B_2(x) &= x \frac{5 - 4x - x^2 + 2\log x + 4x\log x}{2(1-x)^4} \\
P_1(x) &= \frac{1 - 6x + 18x^2 - 10x^3 - 3x^4 + 12x^3\log x}{18(x-1)^5} \\
P_2(x) &= \frac{7 - 18x + 9x^2 + 2x^3 + 3\log x - 9x^2\log x}{9(x-1)^5}
\end{aligned} \tag{42}$$

B_i and P_i are box and penguin diagram contributions respectively.

$$C_{\gamma s}^\pm = -\frac{8\pi Q_d}{3\widetilde{m}^2} (\delta_{12}^d)_{LL} M_4(x), \tag{43}$$

$$C_{\gamma \tilde{g}}^\pm = \frac{8\pi Q_d}{3\widetilde{m}^2} (\delta_{12}^d)_{LL} (\delta_{22}^d)_{LR} (\widetilde{m}\sqrt{x}) M_2(x), \tag{44}$$

$$C_{gs}^\pm = -\frac{2\pi}{\widetilde{m}^2} (\delta_{12}^d)_{LL} \left(\frac{3}{2}M_3(x) - \frac{1}{6}M_4(x) \right), \tag{45}$$

$$C_{g\tilde{g}}^\pm = \frac{2\pi}{\widetilde{m}^2} (\delta_{12}^d)_{LL} (\delta_{22}^d)_{LR} (\widetilde{m}\sqrt{x}) \left(\frac{3}{2}M_1(x) - \frac{1}{6}M_2(x) \right) \tag{46}$$

where $Q_d = -1/3$ and

$$\begin{aligned}
M_1(x) &= \frac{3 - 3x^2 + (1 + 4x + x^2) \ln x}{(x-1)^5} \\
M_2(x) &= \frac{1 + 9x - 9x^2 - x^3 + 6x(1+x) \ln x}{2(x-1)^5} \\
M_3(x) &= \frac{1}{3}M_2(x) \\
M_4(x) &= \frac{-1 + 9x + 9x^2 - 17x^3 + 6x^2(3+x) \ln x}{12(x-1)^5}
\end{aligned} \tag{47}$$

Here again, the Wilson coefficients $C_{\gamma \tilde{g}}^\pm$ and $C_{g\tilde{g}}^\pm$ arising from double mass insertions are given for the first time. These two new terms distinguish our model from other models, when we discuss the processes $K_L \rightarrow \pi^0 e^+ e^-$ and $K_L \rightarrow \mu^+ \mu^-$ through $s \rightarrow d\gamma$, and the process $\text{Re}(\epsilon'/\epsilon_K)$ through $s \rightarrow dg$. The process $K_L \rightarrow \pi^0 e^+ e^-$ is also affected by the $s \rightarrow dl^+ l^-$ local interaction whose Wilson coefficient at the SUSY particle mass scale is

$$C_{7V}^{\text{SUSY}'} = -\frac{4}{9\widetilde{m}^2} P_1(x) (\delta_{12}^d)_{LL}. \tag{48}$$

D. EDM constraint

The Wilson coefficients for the effective Hamiltonian for the neutron EDM in MIA are given by

$$C_1^{\text{edm}} = -\frac{2}{3} \frac{e\alpha_s}{\pi} Q_d \frac{m_{\tilde{g}}}{\widetilde{m}^2} \text{Im}(\delta_{11}^d)_{LR} 4 B_1(x), \quad (49)$$

$$C_2^{\text{edm}} = -\frac{1}{4} \frac{g_s\alpha_s}{\pi} \frac{m_{\tilde{g}}}{\widetilde{m}^2} \text{Im}(\delta_{11}^d)_{LR} \left(\frac{3}{x} B_2(x) - \frac{4}{3} B_1(x) \right), \quad (50)$$

where

$$(\delta_{11}^d)_{LR} = \frac{m_d}{\widetilde{m}^2} (A_d^* - \mu \tan \beta). \quad (51)$$

If we consider the loop functions to be order of one, and set both the squark and the gluino masses to 500 GeV, we get the limit on the flavor preserving mixing parameter,

$$\text{Im}(\delta_{11}^d)_{LR} \lesssim \mathcal{O}(10^{-8}), \quad (52)$$

from the constraint (30). The typical modulus of $(\delta_{11}^d)_{LR}$ is $\mathcal{O}(10^{-4})$ for $A_d^* - \mu \tan \beta = 10\text{TeV}$, and this limit constrains the phase of $(\delta_{11}^d)_{LR}$ below $\mathcal{O}(10^{-4})$.

There is a correlation between phases of $(\delta_{11}^d)_{LR}$ and $(\delta_{22}^d)_{LR}$ because they share a common term $\mu \tan \beta$, and this correlation becomes stronger if we assume universality between A_d and A_s . This indirectly constrains the phase of $(\delta_{22}^d)_{LR}$. But we can safely set its phase to zero while satisfying ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$ at the same time, as we will show in the numerical analysis.

E. Numerical Results

Now in Fig. 1, we show the plots of $\text{Re}(\epsilon'/\epsilon_K)$ as a function of the phase of $(\delta_{12}^d)_{LL}$, where $(\delta_{12}^d)_{LL}$ is parametrized as $re^{i\phi}$. We fixed $A_s^* - \mu \tan \beta = 10\text{ TeV}$. Different values and sign of $A_s^* - \mu \tan \beta$ yield similar results. For a given value of ϕ , r is determined in such a way that ϵ_K becomes the experimental value, and $\text{Re}(\epsilon'/\epsilon_K)$ is computed from those two parameters. As ϕ varies from 180° to 360° for $\gamma = 0^\circ$, the SUSY contribution to $\text{Re}(\epsilon'/\epsilon_K)$ repeats its values in the region of ϕ from 0° to 180° , with the sign flipped. Therefore it is clearly ruled out, and the results are not shown. But as the angle γ varies, the SM contribution either increases or decreases $\text{Re}(\epsilon'/\epsilon_K)$, so that the invisible parts in $\gamma = 0$ graph float up in $\gamma = 60^\circ$ and $\gamma = 120^\circ$ graphs, while some visible parts sink down in $\gamma = 240^\circ$ and $\gamma = 300^\circ$ graphs. Note that the functional relation between r and ϕ governed by the ϵ_K constraint depends on the KM angle γ , and thus plots for nonzero γ are not simple shifts of that for vanishing γ .

Since there may be large contributions to $K^0 - \overline{K}^0$ and $B - \overline{B}^0$ mixing in our model, it is utterly important to measure three angles α, β and γ of the unitarity triangle (UT) in a way independent of these neutral meson mixings, in addition to the more popular ways using $B \rightarrow J/\psi K_S, \pi\pi$, so on. Also, it would be an interesting and important question how much portions of $\text{Re}(\epsilon'/\epsilon_K)$ come from the SM and new physics (SUSY gluino-squark loops in

this work). This question can be answered only if the angle γ is determined by some other methods independent of $B^0 - \bar{B}^0$ mixing, e.g., from $B \rightarrow DK$ or $B \rightarrow \pi K$, etc.. Once the angle γ is known, one can calculate the SM contribution, $\text{Re}(\epsilon'/\epsilon_K)^{\text{SM}}$. Then, the difference between the observed $\text{Re}(\epsilon'/\epsilon_K)$ and the SM would be the new physics contributions.

In our model, $K \rightarrow \pi\nu\bar{\nu}$ decays are not directly affected by the gluino-squark loops. This is because neutrinos couple to neither of photons and gluinos, and thus the gluino-squark loop do not induce $sd\nu\bar{\nu}$ vertex. Still, the shape of the UT, namely the angle γ , can be different from the SM, because of potentially large contributions to ϵ_K (and also to $B^0 - \bar{B}^0$ mixing). Therefore, gluino mediated FCNC will affect the branching ratios for $K \rightarrow \pi\nu\bar{\nu}$ in an indirect way, even if the Wilson coefficients for $s \rightarrow d\nu\bar{\nu}$ remains the same as in the SM. In Fig. 2 (a) and (b), we show the $K_L \rightarrow \pi^0\nu\bar{\nu}$ branching ratio as a function of γ and its correlation with $\text{Re}(\epsilon'/\epsilon_K)$, respectively. There is no definite correlation between $K_L \rightarrow \pi^0\nu\bar{\nu}$ and $\text{Re}(\epsilon'/\epsilon_K)$ in our model, unlike the scenario in which the new physics affects $\text{Re}(\epsilon'/\epsilon_K)$ by a modified sdZ vertex. We note that one cannot make definite predictions for the branching ratio for this decay unless we know the angle γ in the presence of new physics. The branching ratio may be vanishingly small or larger than the SM prediction by a factor of $\lesssim 2$. Therefore, this decay once measured may provide us with invaluable informations on the value of the angle γ , if there is no significant new physics contribution to $s \rightarrow d\nu\bar{\nu}$ upto 4-fold ambiguities, which could be eliminated using information from other rare kaon decays and B meson decays. In Fig. 2 (c) and (d), we show the similar plots for $K^+ \rightarrow \pi^+\nu\bar{\nu}$. Again, the measurement of the $K^+ \rightarrow \pi^+\nu\bar{\nu}$ branching ratio will help to determine γ . And the branching ratio for this decay may be larger than the SM predictions by a factor of ~ 2 .

In Figs. 3 (a) and (b), we show the branching ratio of $B(K_L \rightarrow \pi^0 e^+ e^-)_{\text{dir}}$ as a function of ϕ , the argument of $(\delta_{12}^d)_{LL}$, and its nontrivial correlation with $\text{Re}(\epsilon'/\epsilon_K)$. This nontrivial correlation arises because electrons couple to a photon and also the $s \rightarrow d\gamma$ vertex contributes both $K_L \rightarrow \pi^0 e^+ e^-$ and $\text{Re}(\epsilon'/\epsilon_K)$. In other words, the Wilson coefficient for $s \rightarrow d\gamma$ is modified by the gluino mediated FCNC, which is in turn closely related with the gluino mediated sdg operator. Also the Wilson coefficient of O_{7V} vertex is modified by the gluino mediated FCNC, which affects the direct CP violating part in $K_L \rightarrow \pi^0 e^+ e^-$. This is in contrast with the $K \rightarrow \pi\nu\bar{\nu}$ case which is not affected directly by gluino mediated FCNC. Therefore, the direct CP violating part of $K_L \rightarrow \pi^0 e^+ e^-$ can be affected a lot, if there is a new CP violating phase in $\tilde{s}_L - \tilde{d}_L$ mixing both by a single MI of $(\delta_{12}^d)_{LL}$ and a double MI of $(\delta_{12}^d)_{LR}^{\text{ind}} = (\delta_{12}^d)_{LL} \times (\delta_{22}^d)_{LR}$. Also this process is affected indirectly by SUSY gluino effects through their dependence on the KM angle γ . After all, the branching ratio of $K_L \rightarrow \pi^0 e^+ e^-$ can be vanishingly small or enhanced over the SM value by a factor of ~ 3 .

Finally, we consider the gluino mediated SUSY contributions to the short distance part of $K_L \rightarrow \mu^+ \mu^-$. Again this decay could be affected by gluino mediated FCNC both directly and indirectly : through the modified Wilson coefficient of $sd\gamma$ operator and modified values of γ due to SUSY contributions to ϵ_K . However, unlike the direct CP violating part in $K_L \rightarrow \pi^0 e^+ e^-$, the modified $sd\gamma$ is irrelevant to $K_L \rightarrow \mu^+ \mu^-$ because of current conservation associated with the muon current. Therefore it depends only on the angle γ [Fig. 4 (a)], and its correlation with ϵ'/ϵ is rather trivial [Fig. 4 (b)]. Depending on the values of γ , the result can be reduced by a factor of ~ 3 or increased by a factor of ~ 2 .

Let us summarize the results on rare kaon decays within mass insertion approximation. We assumed there are significant gluino-squark loop contributions to ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$

through $(\delta_{12}^d)_{LL} \sim O(10^{-2} - 10^{-3})$ and the $(\delta_{12}^d)_{LR} \sim O(10^{-5})$ induced by a double mass insertion of flavor preserving $\tilde{s}_R - \tilde{s}_L$ followed by $(\delta_{12}^d)_{LL}$. Then these parameters directly affect $s \rightarrow dg$ and $s \rightarrow d\gamma$ vertices, but not $s \rightarrow d\nu\bar{\nu}$. Also the shape of the UT would be different from the SM case because SUSY loop contributes to ϵ_K . Therefore all the rare kaon decays (and also B decays as well, although we do not discuss this subject in any detail here) are affected indirectly through modified CKM elements, especially γ . For $K \rightarrow \pi\nu\bar{\nu}$ and the short distance part of $K_L \rightarrow \mu^+\mu^-$, the branching ratios can be larger than the SM expectations by a factor of $\sim 2 - 3$, or can be smaller than the SM values. Especially the branching ratio for $K_L \rightarrow \pi^0\nu\bar{\nu}$ can be vanishingly small if $\gamma = 0$ or π . Their correlations with $\text{Re}(\epsilon'/\epsilon_K)$ are rather trivial, because these decays are not affected directly by $(\delta_{12}^d)_{LL}$. On the other hand, the direct CP violating part of $K_L \rightarrow \pi^0 e^+ e^-$ is affected both directly and indirectly, and its correlation with $\text{Re}(\epsilon'/\epsilon_K)$ is nontrivial at all. One may expect much enhanced (or reduced) direct CP violations in $K_L \rightarrow \pi^0 e^+ e^-$ in our scenario, depending whether SUSY effects can interfere constructively (destructively) with the SM amplitude.

IV. VERTEX MIXING CASE

A. Flavor mixings in the quark-squark-gluino vertices

Another way of solving the SUSY flavor and CP problem is to assume that the 1st/2nd generation squarks are heavy and there are certain degrees of hierarchy in the $q_i - \tilde{q}_j - \tilde{g}$ mixing matrices, W_L and W_R (the decoupling scenario). Here W_L and W_R are analogous to the CKM matrix in the SM. Interactions among quarks, squarks, and gluinos in terms of their mass eigenstates are described by the Lagrangian:²

$$\mathcal{L} = -\sqrt{2} g_s (W_L)_{ij} \tilde{q}_L^{*i} \tilde{g}_R^a t^a q_L^j + \sqrt{2} g_s (W_R)_{ij} \tilde{q}_R^{*i} \tilde{g}_L^a t^a q_R^j + \text{h.c.}, \quad (53)$$

where i and j are generation indices and color indices of (s)quarks have been omitted. Although squarks of each chirality are diagonalized into their (approximate) mass eigenstates, the remaining small mixing between squarks of different chiralities can be treated as a perturbation except for the stop sector.

B. Wilson coefficient for $\Delta S = 2$

The Wilson coefficients for $\Delta S = 2$ effective Hamiltonian in the VM case are given by following expressions :

$$C_1^{(2)} = +\frac{1}{4} \left(24f_4(x) + 66\tilde{f}_4(x) \right) (F_{12})_{LL}^2,$$

$$C_2^{(2)} = +\frac{204}{4} f_4(x) (F_{12})_{RL}^2,$$

²Our definition of the mixing matrix W is different from that in Ref. [46] in that the signs of the two terms are opposite and both W_L and W_R are not complex-conjugated.

$$\begin{aligned}
C_3^{(2)} &= -\frac{36}{4}f_4(x)(F_{12})_{RL}^2, \\
C_4^{(2)} &= +\frac{1}{4}\left(504\tilde{f}_4(x) - 72f_4(x)\right)(F_{12})_{LL}(F_{12})_{RR} - \frac{132}{4}\tilde{f}_4(x)(F_{12})_{LR}(F_{12})_{RL}, \\
C_5^{(2)} &= +\frac{1}{4}\left(24\tilde{f}_4(x) + 120f_4(x)\right)(F_{12})_{LL}(F_{12})_{RR} - \frac{180}{4}\tilde{f}_4(x)(F_{12})_{LR}(F_{12})_{RL}, \\
\tilde{C}_1^{(2)} &= +\frac{1}{4}\left(24f_4(x) + 66\tilde{f}_4(x)\right)(F_{12})_{RR}^2, \\
\tilde{C}_2^{(2)} &= +\frac{204}{4}f_4(x)(F_{12})_{LR}^2, \\
\tilde{C}_3^{(2)} &= -\frac{36}{4}f_4(x)(F_{12})_{LR}^2,
\end{aligned} \tag{54}$$

where

$$\begin{aligned}
f_4(x) &= x \frac{2x - 2 - (1+x)\ln x}{(x-1)^3}, \\
\tilde{f}_4(x) &= \frac{x^2 - 1 - 2x\ln x}{(x-1)^3}, \\
(F_{12})_{AB} &= (W_A)_{31}^*(W_B)_{32}. \quad (A, B = L \text{ or } R)
\end{aligned} \tag{55}$$

Here we assume that both left and right sbottom masses are equal to \tilde{m} , and define $x = (m_{\tilde{g}}/\tilde{m})^2$.

C. Wilson coefficient for $\Delta S = 1$

The Wilson coefficients for $\Delta S = 1$ effective Hamiltonian in the VM case are given by following expressions :

$$C_i = 0 \quad (\text{for } i = 1, 2, 7 \sim 10),$$

$$C_3 = \frac{1}{4} \frac{\alpha_s^2}{2\tilde{m}^2} (F_{12})_{LL} P(x),$$

$$C_4 = -\frac{1}{4} \frac{3\alpha_s^2}{2\tilde{m}^2} (F_{12})_{LL} P(x),$$

$$C_5 = \frac{1}{4} \frac{\alpha_s^2}{2\tilde{m}^2} (F_{12})_{LL} P(x),$$

$$C_6 = -\frac{1}{4} \frac{3\alpha_s^2}{2\tilde{m}^2} (F_{12})_{LL} P(x),$$

$$\tilde{C}_i = C_i \text{ with replacement of } LL \text{ by } RR, \quad i = 3, 4, 5, 6$$

$$C_{\gamma s}^{\pm} = \frac{8\pi Q_d}{3\tilde{m}^2} [(F_{12})_{LL} \pm (F_{12})_{RR}] P_{BE}(x), \tag{56}$$

$$C_{\gamma \tilde{g}}^{\pm} = \frac{8\pi Q_d}{3\tilde{m}^2} \left[(F_{12})_{LR} \frac{m_b(A_b^* - \mu \tan \beta)}{\tilde{m}} \pm (F_{12})_{RL} \frac{m_b(A_b - \mu^* \tan \beta)}{\tilde{m}} \right] \sqrt{x} P_{BI}(x, x), \tag{57}$$

$$C_{gs}^{\pm} = \frac{2\pi}{\tilde{m}^2} [(F_{12})_{LL} \pm (F_{12})_{RR}] P_2(x), \tag{58}$$

$$C_{g\tilde{g}}^{\pm} = \frac{2\pi}{\tilde{m}^2} \left[(F_{12})_{LR} \frac{m_b(A_b^* - \mu \tan \beta)}{\tilde{m}} \pm (F_{12})_{RL} \frac{m_b(A_b - \mu^* \tan \beta)}{\tilde{m}} \right] \sqrt{x} P_3(x, x), \tag{59}$$

where the LR mixing parameter $(A_b^* - \mu \tan \beta)$ is real, and

$$\begin{aligned}
P_F(x) &= \frac{1}{36} \frac{1}{(1-x)^4} \left[7x^3 - 36x^2 + 45x - 16 + (18x - 12) \ln x \right], \\
P_B(x) &= \frac{-1}{36} \frac{1}{(1-x)^4} \left[11x^3 - 18x^2 + 9x - 2 - 6x^3 \ln x \right], \\
P_{FE}(x) &= \frac{1}{12} \frac{-1}{(1-x)^4} \left[x^3 - 6x^2 + 3x + 2 + 6x \ln x \right], \\
P_{BE}(x) &= \frac{1}{12} \frac{1}{(1-x)^4} \left[2x^3 + 3x^2 - 6x + 1 - 6x^2 \ln x \right], \\
P_{FI}(x) &= \frac{1}{2} \frac{-x}{(1-x)^3} \left[4x - x^2 - 3 - 2 \ln x \right], \\
P_{BI}(x) &= \frac{1}{2} \frac{x}{(1-x)^3} \left[1 - x^2 + 2x \ln x \right], \\
P(x) &= P_F(x) - \frac{1}{9} P_B(x), \\
P_2(x) &= \frac{3}{2} P_{FE}(x) - \frac{1}{6} P_{BE}(x), \\
P_3(x) &= \frac{3}{2} P_{FI}(x) - \frac{1}{6} P_{BI}(x), \\
P_{BI}(x, x) &= \lim_{y \rightarrow x} \frac{P_{BI}(x) - P_{BI}(y)}{x - y}, \\
P_3(x, x) &= \lim_{y \rightarrow x} \frac{P_3(x) - P_3(y)}{x - y}.
\end{aligned} \tag{60}$$

We don't consider the box diagrams in vertex mixing, because it will be highly suppressed due to heavy 1st/2nd generation squarks.

The Wilson coefficient of $O_{7V}^{\text{SUSY}'}$ is

$$C_{7V}^{\text{SUSY}'} = -\frac{4}{9\tilde{m}^2} P_B(x) (F_{12})_{LL}. \tag{61}$$

D. EDM constraint

The Wilson coefficients for the effective Hamiltonian for the neutron EDM in VM are given by

$$\begin{aligned}
C_1^{\text{edm}} &= -\frac{2}{3} \frac{e\alpha_s}{\pi} Q_d \frac{m_{\tilde{g}}}{\tilde{m}^2} \frac{m_b \text{Im}[(F_{11})_{LR}(A_b^* - \mu \tan \beta)]}{\tilde{m}^2} P_{BI}(x, x), \\
C_2^{\text{edm}} &= -\frac{1}{4} \frac{g_s\alpha_s}{\pi} \frac{m_{\tilde{g}}}{\tilde{m}^2} \frac{m_b \text{Im}[(F_{11})_{LR}(A_b^* - \mu \tan \beta)]}{\tilde{m}^2} 2 \left(\frac{3}{2} P_{FI} - \frac{1}{6} P_{BI} \right) (x, x).
\end{aligned} \tag{62}$$

Basically, the constraint on $\text{Im}(\delta_{11}^d)_{LR}$ in (52), applies to $m_b \text{Im}[(F_{11})_{LR}(A_b^* - \mu \tan \beta)]/\tilde{m}^2$ much the same way, and if we assume $(A_b^* - \mu \tan \beta)$ is real and is about 2 TeV, this limits

$$\text{Im}(F_{11})_{LR} \lesssim \mathcal{O}(10^{-7}). \tag{63}$$

Since we will set $(W_R)_{31}$ to zero in the numerical analysis, there is no difficulty in satisfying this limit in our case, but it should be obeyed when other possible solutions are searched for. For example, if both $(W_L)_{31}$ and $(W_R)_{31}$ have absolute values of $\mathcal{O}(10^{-2})$, the typical order of magnitude of $|(W_L)_{31}|$ in our analysis, then their phase difference should be less than $\mathcal{O}(10^{-3})$.

E. Numerical Results

In MI case, the only key parameter responsible for CP violation was $(\delta_{12}^d)_{LL}$. Once its modulus is determined by ϵ_K constraint for a fixed γ , every CP violating observable could be computed as a function of its phase, $\text{Arg}(\delta_{12}^d)_{LL}$. Within the VM approach, on the contrary, the analysis becomes more complicated, because we have four independent complex parameters, $(W_L)_{31}$, $(W_L)_{32}$, $(W_R)_{31}$, and $(W_R)_{32}$, that describe the flavor mixing in the $q - \tilde{q} - \tilde{g}$ with various chiralities of quarks and their superpartners. Even if we impose the ϵ_K and the EDM constraints, there still remain 6 real parameters and it would be a formidable task to scan all of this 6 dimensional parameter space. In order to simplify the numerical analysis, we discarded some parameters and tried to find a particular solution rather than look for a completely general solution. That would not much distort the picture of VM approach and would suffice to illustrate a possible existence of appropriate parameter space in which gluino-squark loop contributions to ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$ are significant without conflict with electron/neutron EDM's.

For definiteness, we set $(W_L)_{32}$ and $(W_R)_{31}$ to zero, and all SUSY contributions are expressed in terms of a single complex number $(F_{12})_{LR} \equiv (W_L)_{31}^* (W_R)_{32}$. The results for $\text{Re}(\epsilon'/\epsilon_K)$ for various values of the angle γ 's are displayed in Figs. 5, as functions of ϕ , where we use the following parameterization, $(F_{12})_{LR} = r e^{i\phi}$. As in the MIA case, we can have a large ϵ'/ϵ in certain region of ϕ with size $\sim \mathcal{O}(1)$. In particular, all the CP violations in the kaon system could be accommodated in terms of CP violations from SUSY sector with vanishing KM angle, $\gamma = 0$ or π , if $A_b^* - \mu \tan \beta = \mathcal{O}(2)$ TeV. Therefore, $\tan \beta$ could be considerably smaller than the MIA case because the flavor preserving LR mixing in the VM case is enhanced by m_b/m_s relative to that in the MIA. Also SUSY particle mass spectra will be quite different from those in the MIA case, although both cases can accommodate the CP violations in the kaon system.

The decays $K \rightarrow \pi \nu \bar{\nu}$ and the short distance part of $K_L \rightarrow \mu^+ \mu^-$ in the VM will be essentially the same as those in the MIA. There are no direct SUSY effects on these processes. SUSY effects are only indirect through the angle γ from the SUSY contributions to ϵ_K . Therefore we do not show any plots for these processes. The γ dependence are the same as the MIA [see Figs. 2 (a), (c) and Fig. 4 (a)].

On the contrary, the direct CP violating part of $K_L \rightarrow \pi^0 e^+ e^-$ is affected directly by the SUSY loop contributions. In Fig. 6 (a), we show its branching ratio as a function of the angle ϕ for different values of γ . In Fig. 6 (b), we show its correlation with $\text{Re}(\epsilon'/\epsilon_K)$. It looks very similar to Fig. 3 (b) because $C_\gamma^+ = C_\gamma^-$ and $C_g^+ = C_g^-$ here as was in the MIA case. If we relax the assumption that $(W_L)_{32} = (W_R)_{31} = 0$, however, then the correlation may differ from that in the MIA case in general because $\text{Re}(\epsilon'/\epsilon_K)$ mainly depends on C_g^- while $B(K_L \rightarrow \pi^0 e^+ e^-)$ on C_γ^+ .

V. CONCLUSIONS

In this paper, we discussed the gluino–squark loop contributions to CP violations in the kaon system both in the MIA and in the VM. Unlike many other models based on the enhanced sdg or sdZ vertex in the context of SUSY using nonuniversal flavor structures in the trilinear A couplings, our model modifies both the ε and $\text{Re}(\epsilon'/\epsilon_K)$, so that the phenomenological implications on other kaon decays substantially differ from other models. In particular the shape of unitarity triangle can be very different from the SM case, due to potentially large SUSY contributions to ϵ_K and $B^0 - \overline{B}^0$ mixing. Only the constraint from the semileptonic $b \rightarrow u$ transition is stable against a possible SUSY contributions. Therefore we varied the γ from 0 to 2π , then fixed the λ_t to get the right magnitude of ε_K . In particular, one can accommodate both ϵ_K and $\text{Re}(\epsilon'/\epsilon_K)$, even if the angle γ is very close to 0 or π , still consistent with the ranges obtained from the three generation unitarity relations, Eq. (2). In our model, the gluino–squark loop does modify $s \rightarrow d\gamma$ just like the sdg operator, but does not directly contribute to $s \rightarrow d\nu\bar{\nu}$. Still the branching ratios for $K \rightarrow \pi\nu\bar{\nu}$ can be different from the SM predictions by a large amount because of the modified γ . Similar effects occur in $K_L \rightarrow \pi^0 e^+ e^-$ and $K_L \rightarrow \mu^+ \mu^-$ processes. If the shape of the triangle in MSSM is the same as in the SM case, then the deviations from the SM predictions are possible only in ϵ'/ϵ_K and $K_L \rightarrow \pi^0 e^+ e^-$, although the deviation in the latter case is very small. Also in our model, SUSY effects give additional contributions to $\Delta S = 1$ penguin operators, so that we may anticipate interesting deviations of direct CP asymmetries in $K \rightarrow 3\pi$ and $K \rightarrow \pi\pi\gamma$ from the SM predictions. However, these observables are generically contaminated by hadronic uncertainties even in the chiral perturbation theory, and are left for the future study.

One can also perform a similar analysis in the B meson sector including SUSY gluino–squark loop contributions to $B^0 - \overline{B}^0$ mixing, and considering a specific flavor structure for the mass insertion parameters [47]. Depending on the chiral structures of the mass insertion parameters, one could achieve $\sin 2\beta_{J/\psi K_s}$ which is very different from the SM case. In this case, the results could be very different from the SM even if we assume that the UT in the MSSM is the same as in the SM.

Our study indicates that the measurements of $K_L \rightarrow \pi^0 \nu\bar{\nu}$ in the kaon sector is very important for constructing the unitarity triangle, especially if there are large new physics contributions to ϵ_K and $B^0 - \overline{B}^0$ mixing. Also, in the B meson sector, it would be utterly important to construct the UT from the tree level processes which are immune to possible new physics contributions. For example, extracting γ from $B \rightarrow DK$ or other tree dominated decays will be very desirable [48]. The current global analysis on CKM matrix elements suggests that a unique UT emerges and the SM picture for flavor and CP violations in the quark sector seems to be a correct picture. But as we demonstrated with our specific model, this is no longer guaranteed in the presence of new physics which give significant contributions to ϵ_K and $B^0 - \overline{B}^0$ mixing [49].

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FIGURES

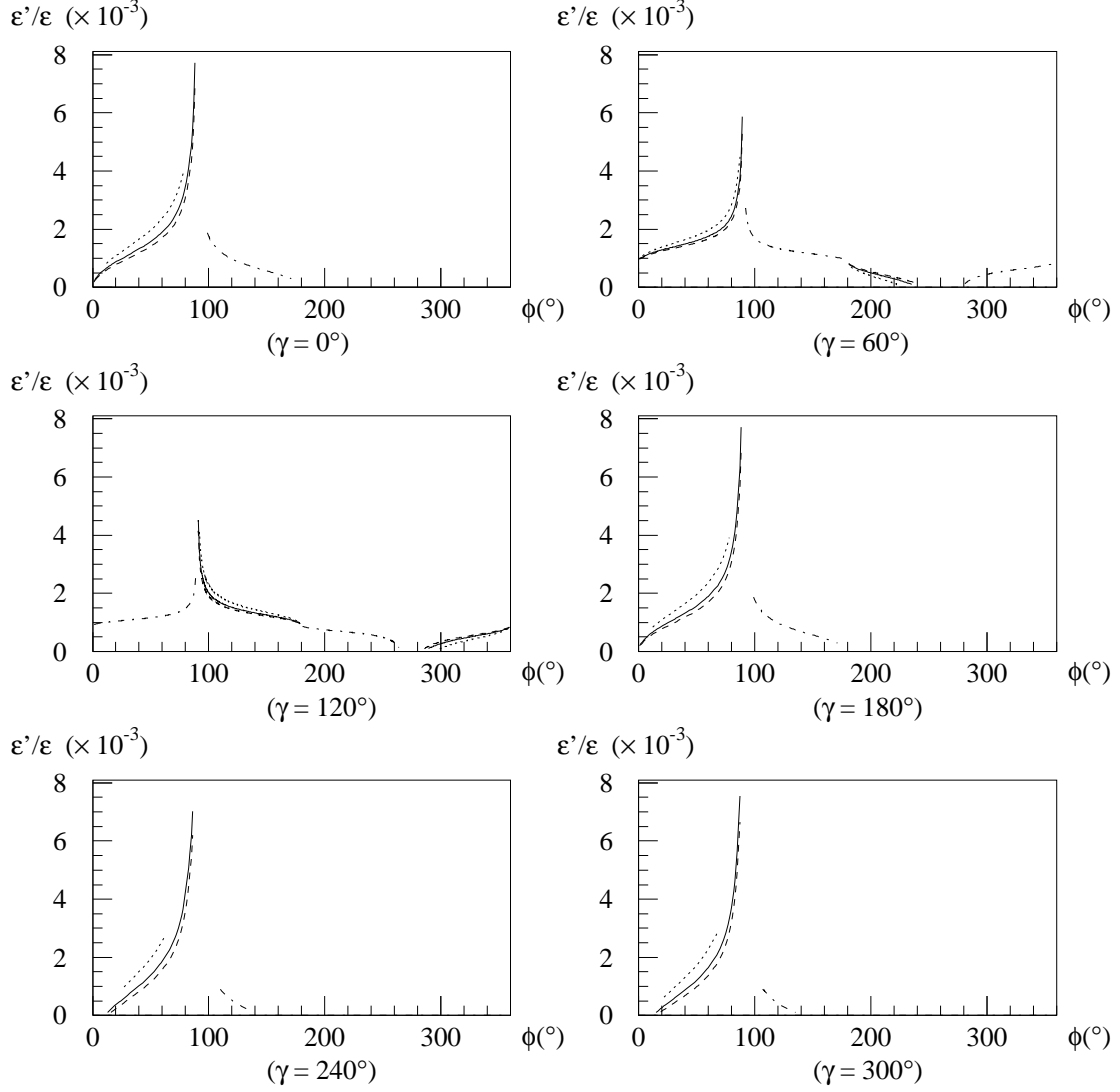


FIG. 1. The plots of ϵ'/ϵ versus the phase of $(\delta_{12}^d)_{LL}$ in MI for six different values of γ with $A_s^* - \mu \tan \beta = 10 \text{ TeV}$. Graphs were drawn in the solid, the dashed, the dotted, and the dash-dotted lines for $x = 0.3, 1.0, 2.0, 4.0$, respectively.

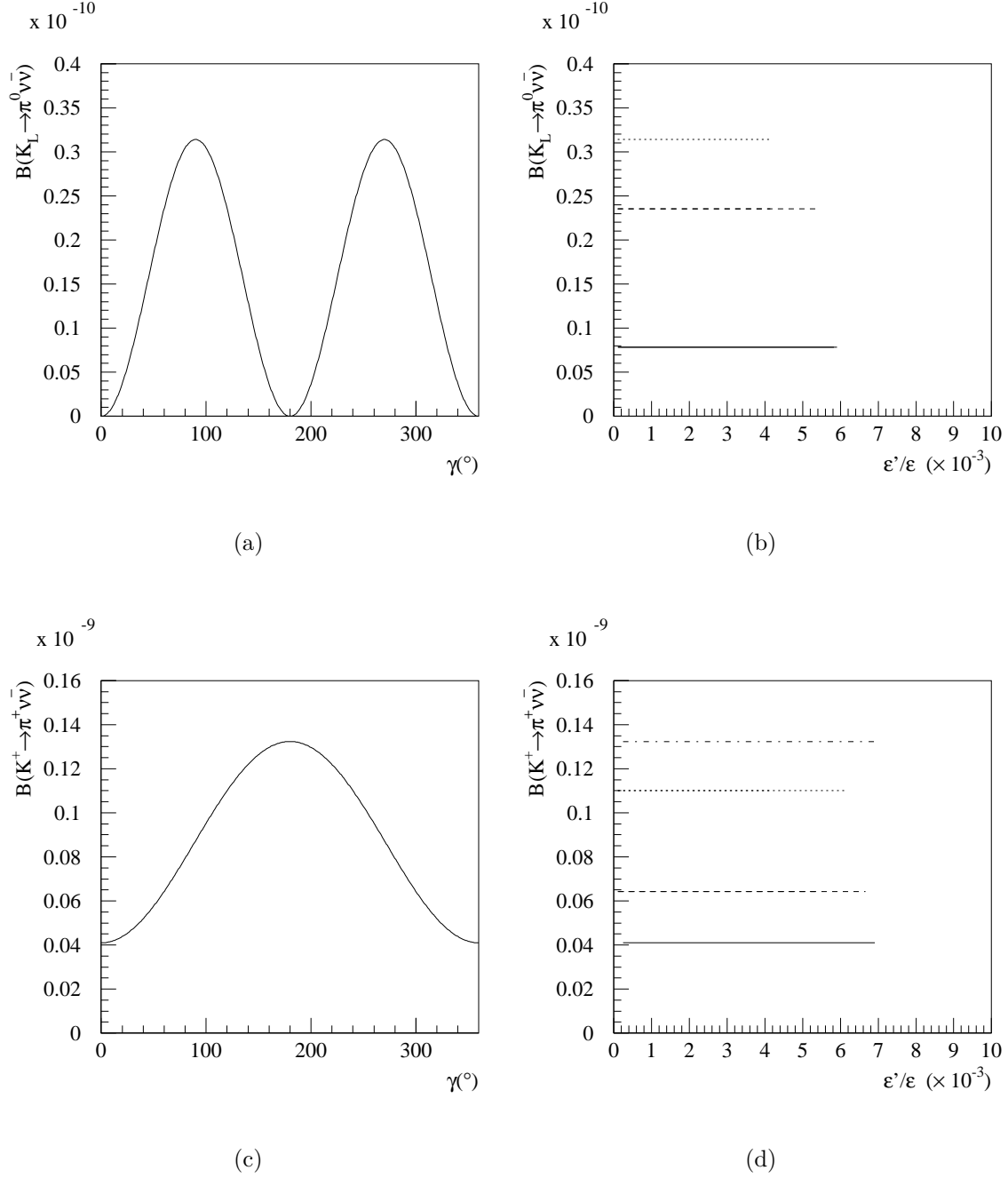


FIG. 2. The γ dependences of $B(K \rightarrow \pi \nu \bar{\nu})$ ((a), (c)) and their correlations with ϵ'/ϵ in MI ((b), (d)), for $x = 1$ with $A_s^* - \mu \tan \beta = 10$ TeV. (b) The solid, the dashed, and the dotted lines correspond to $\gamma = 30^\circ, 60^\circ, 90^\circ$, respectively. The first two equally correspond to $\gamma = 150^\circ, 120^\circ$ as well. The branching ratio vanishes for $\gamma = 0^\circ$ or 180° . (d) The solid, the dashed, the dotted, and the dash-dotted lines correspond to $\gamma = 0^\circ, 60^\circ, 120^\circ, 180^\circ$, respectively, and 360° minus them.

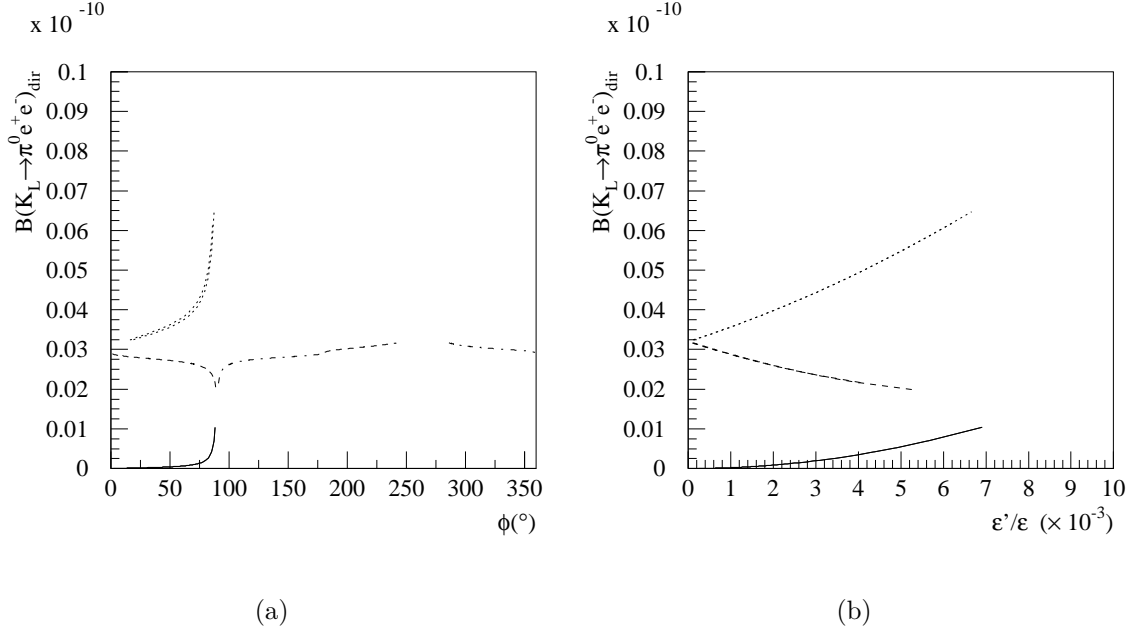


FIG. 3. The ϕ dependence of $B(K_L \rightarrow \pi^0 e^+ e^-)_{\text{dir}}$ ((a)) and its correlation with ϵ'/ϵ ((b)), in MI for $x = 1$ with $A_s^* - \mu \tan \beta = 10$ TeV. (a) The solid, the dashed, the dotted, and the dash-dotted lines correspond to $\gamma = 0^{\circ}, 180^{\circ}; 60^{\circ}; 240^{\circ}, 300^{\circ}; 120^{\circ}$, respectively. (b) The solid, the dashed, and the dotted lines correspond to $\gamma = 0^{\circ}, 60^{\circ}, 240^{\circ}$, respectively, and 180° minus them.

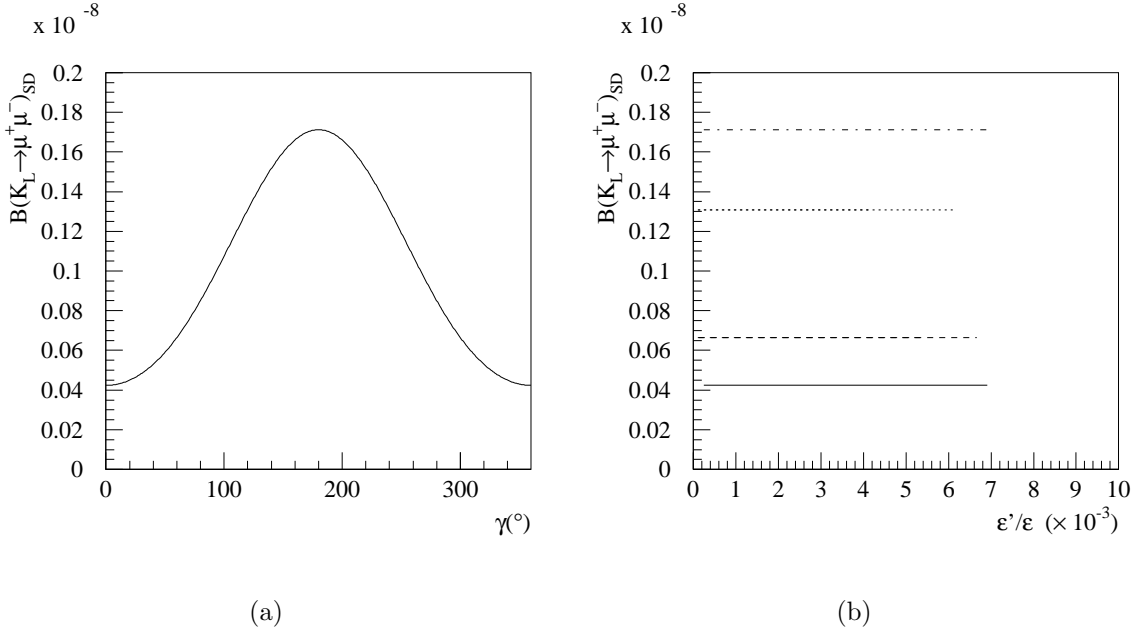


FIG. 4. The γ dependence of $B(K_L \rightarrow \mu^+ \mu^-)_{\text{SD}}$ ((a)) and its correlation with ϵ'/ϵ in MI ((b)), for $x = 1$ with $A_s^* - \mu \tan \beta = 10$ TeV. (b) The solid, the dashed, the dotted, and the dash-dotted lines correspond to $\gamma = 0^{\circ}, 60^{\circ}, 240^{\circ}, 180^{\circ}$, respectively, and 360° minus them.

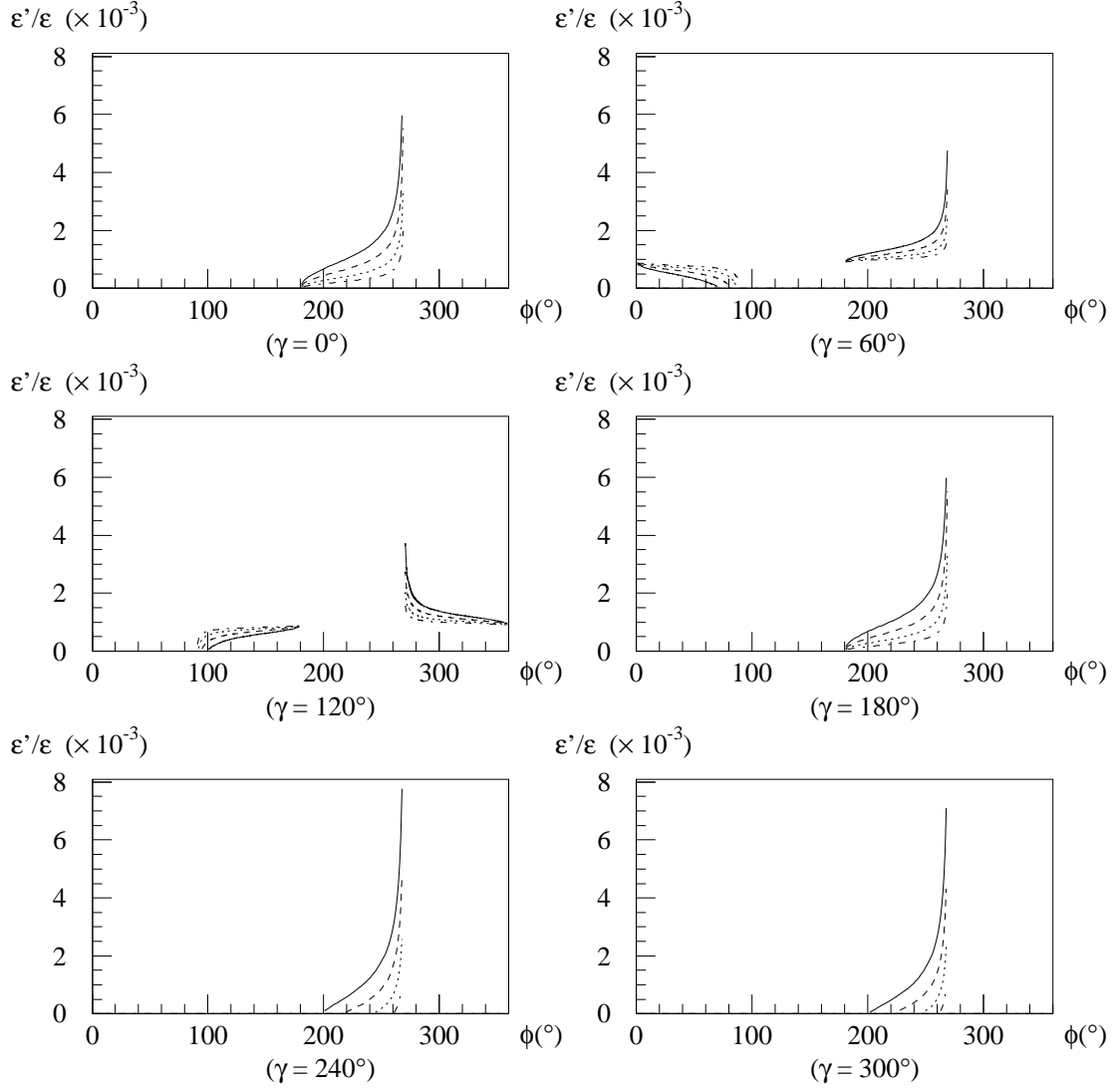
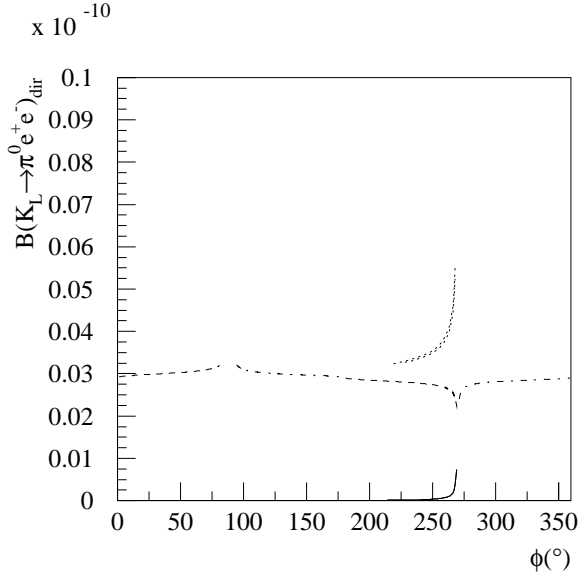
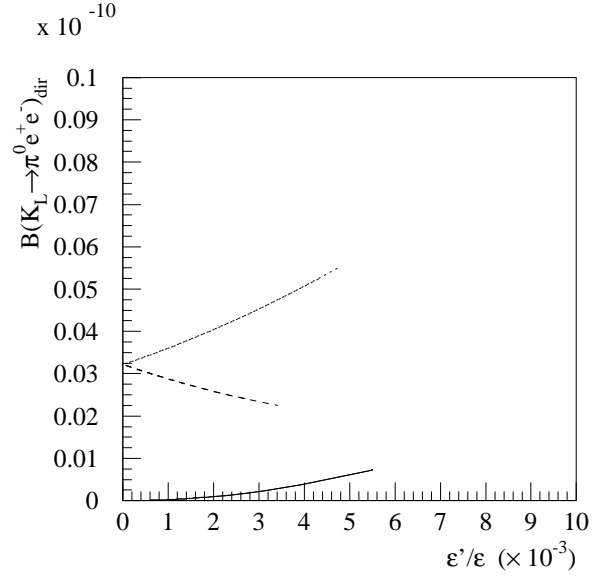


FIG. 5. The plots of ϵ'/ϵ versus the phase of $(F_{12})_{LR}$ in VM for six different values of γ with $A_b^* - \mu \tan \beta = 2\text{TeV}$. Graphs were drawn in the solid, the dashed, the dotted, and the dash-dotted lines for $x = 0.5, 1.0, 2.0, 4.0$, respectively.



(a)



(b)

FIG. 6. The ϕ dependence of $B(K_L \rightarrow \pi^0 e^+ e^-)_{\text{dir}}$ ((a)) and its correlation with ϵ'/ϵ ((b)), in VM for $x = 1$ with $A_s^* - \mu \tan \beta = 2$ TeV. These curves are specific to the solution of parameter space drawn in Fig. 5 which restricts $(W_L)_{32}$ and $(W_R)_{31}$ to zero. The meanings of line patterns are the same as in Fig. 3.

TABLES

m_Z	91.2 GeV	$\alpha_s(m_Z)$	0.118
m_W	80.2 GeV	f_π	131 MeV
m_t	170 GeV	f_K	160 MeV
m_b	4.4 GeV	ΔM_K	3.51×10^{-15} GeV
m_c	1.3 GeV	$\sin^2 \theta_W$	0.23
m_d	8 MeV	$ \epsilon_K $	$(2.266 \pm 0.023) \times 10^{-3}$
m_π	135 MeV	$(\text{Re}A_0)_{\text{exp}}$	3.33×10^{-7}
m_K	498 MeV	$(\text{Re}A_2)_{\text{exp}}$	1.50×10^{-8}

TABLE I. Input values we used in the numerical analysis.